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Gauge mediated supersymmetry breaking and supergravity

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Abstract

We analyze simple models of gauge mediated supersymmetry breaking in the context of supergravity. We distinguish two cases. One is when the messenger of the supersymmetry breaking is a non Abelian gauge force and the other is when the messenger is a pseudoanomalous $U(1)$. We assume that these models originate from string theory and we impose the constraint of the vanishing of the cosmological constant. In the first case, we do not find vacua that are consistent with the constraints of gauge mediation and have a zero tree level cosmological constant. In the second case, no such conflict arises. In addition, by looking at the one loop cosmological constant, we show that the dilaton F-term can not be neglected in either case.

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1 Introduction

One possible scenario of supersymmetry breaking is when there is a hidden sector in which gaugino condensation takes place [1]. The breakdown of supersymmetry can be communicated to the visible sector either by gravity or by a non gravitational gauge interaction. In the second class of models we can distinguish two possibilities. One, is when the information of supersymmetry breaking is communicated to the visible sector by the usual gauge (non-Abelian) interactions [2]. This mechanism is called Gauge Mediated Supersymmetry Breaking (GMSB), an old idea recently revived in [3]. The quantity that sets the scale for the soft masses is $\frac{\alpha}{4\pi}\frac{F_X}{<X>}$, where F_X and $< X >$ are the highest and lowest components of a chiral superfield X in the hidden sector. Phenomenological reasons [3] require $< X >$ to be rather low compared to the cut-off scale $M \sim M_{Pl}$. The other, is when supersymmetry breaking is communicated to the visible sector by a pseudoanomalous $U(1)$ gauge interaction (U1MSB) [5]. The scale of the soft masses in this case is set either by the D-term corresponding to the $U(1)$ or, in the absence of D-term contributions, it is set by an F-term. Since the virtue of the pseudoanomalous $U(1)$ is mainly the generation of fermion masses, it is assumed to be flavor non universal over the visible sector. Then, however, in order to avoid conflict with data on flavor changing neutral currents (fcnc), we have to require that the D-terms essentially do not contribute to supersymmetry breaking. Furthermore, to have universal masses at the scale M, (again for fcnc reasons) we would like, in addition, supersymmetry breaking to be dominated by the dilaton F-term, so that the superpartners would get universal masses $\sim F_S/M$. We will analyze prototype models for each case and we will try to answer the question under what circumstances supersymmetry is unbroken or broken with vanishing cosmological constant and what are the main features of the corresponding unbroken or broken vacua.

For GMSB, we will take the tree level superpotential to be simply $w_0 = \lambda X\Phi\bar{\Phi}$, with λ a Yukawa coupling. Φ and $\bar{\Phi}$ are the messenger fields, vector-like with respect to the gauge

[3] One of the constraints on $< X >$ comes from standard Big Bang nucleosynthesis, giving an upper bound of $< X > \leq 10^{12}\ GeV$ [4].

group that contains the standard model gauge group. X is a field, singlet of the visible and hidden sector gauge groups. For concreteness, we will take a specific, $SU(N_c)$ hidden sector that contains one family ($N_f = 1$) of fields Z and $\bar{Z}$ transforming respectively as $\mathbf{N}$ and $\bar{\mathbf{N}}$ of $SU(N_c)$. If the hidden sector is asymptotically free, then at the confining scale the formation of gaugino condensates becomes possible. Below this scale, the relevant physical hidden sector field becomes the hidden sector singlet "meson" $X \equiv \sqrt{Z\bar{Z}}$, and the corresponding contribution to the superpotential, for $N_c > N_f + 1$, will be [6] $w(S,X) = c \cdot X^{-p} e^{-rS}$. Here c, p and r are model dependent parameters and S is the dilaton field. Given that we assumed one hidden sector flavor, $p = 2/(N_c - N_f)$ and $r = 8\pi^2/(N_c - N_f)$ are both strictly positive quantities. This specific choice of the hidden sector, will not spoil the generality of our final conclusions concerning supersymmetry breaking and the vanishing of the cosmological constant. The full superpotential is then $\mathcal{W} = w_0(X, \Phi, \bar{\Phi}) + w(S,X) + k$. In $\mathcal{W}$ we allow for a constant term k which can be either the vacuum value of a combination of fields that have been set to constants upon minimization of the scalar potential with respect to the corresponding fields, or a new, non perturbative contribution to $\mathcal{W}$. In either case what is important for us is that k is independent of S, X, Φ and $\bar{\Phi}$.

For U1MSB, we assume a hidden sector which has exactly the same structure as the hidden sector of the GMSB model and we couple it to a visible sector singlet field Θ. Below the condensation scale, the physical degree of freedom is $\sqrt{Z\bar{Z}} \equiv X$. Both Θ and X are assumed to be charged under the $U(1)$. The full superpotential in this case is taken to be $\mathcal{W} = w_0 + w + k = \lambda \Theta X^2 + c \cdot X^{-p} e^{-rS} + k$, where k is independent of Θ and X

Due to the dilaton dependence of w, it is naturally implied that these superpotentials arise from some string compactification which requires supergravity as the correct low energy effective theory. In supergravity, the scalar potential (ignoring D-terms), is given by [7]:

$$V = e^{K/M^2} \Big[\sum_{\phi_i} |F_{\phi_i}|^2 - 3 \frac{|\mathcal{W}|^2}{M^6} \Big] \quad \text{with} \quad F_{\phi_i} \equiv \mathcal{W}^{(\phi_i)} + \frac{K^{(\phi_i)}}{M^2} \mathcal{W}, \tag{1.1}$$

where K is the Kähler potential, and ϕ_i denotes any of the physical fields. Superscripts in parentheses denote differentiation with respect to the corresponding fields. We will assume for simplicity that all the other moduli besides the dilaton (such as T or U) can be neglected. The Kähler potential K then will simply be:

$$K = -M^2 \log(S + \bar{S}) - \sum_{\phi_i} \phi_i \phi_i^* . \tag{1.2}$$

It is also convenient to set $M = 1$ and to define $\mathcal{G} \equiv K - \log|\mathcal{W}|^2$. The task is to minimize the potential (1.1) for the GMSB and U1MSB models and see if supersymmetry breaking vacua with desired phenomenological properties exist.

In section 2, we investigate supersymmetric vacuum configurations with vanishing tree level cosmological constant. A necessary and sufficient condition for this is:

$$F_{\phi_i} = 0\ , \quad D = 0 \quad \text{and} \quad \mathcal{W} = 0. \tag{1.3}$$

The $D = 0$ condition is necessary in the $U(1)$ case. In fact, we will show that there are no supersymmetric vacua for either case with a nonzero gauge coupling. In section 3, we investigate supersymmetry breaking. To look for supersymmetry breaking vacua, we have to minimize the potential. The scalar potential $V = e^K[\cdots]$ upon minimization, gives $V^{(\phi_i)} = K^{(\phi_i)} e^K[\cdots] + e^K[\cdots]^{(\phi_i)} = 0$, which implies that in a vacuum with vanishing tree level cosmological constant it is sufficient to solve for $[\cdots]^{(\phi_i)} = 0$, provided that $K \neq -\infty$ at minimum. In section 4, we discuss vanishing of the cosmological constant at one loop. In section 5, we give our conclusions.

2 Supersymmetric Vacua

GMSB model

Let us start by writing out our superpotential again as [2]:

$$\mathcal{W} = w_0 + w + k = \lambda X \Phi \bar{\Phi} + c X^{-p} e^{-rS} + k. \tag{2.4}$$

Using the Kähler potential (1.2), the F-terms are computed to be:

$$F_X = \frac{1}{X}(w_0 - pw) - X^* \mathcal{W}\ , \quad F_S = -rw - \sigma \mathcal{W}, \tag{2.5}$$

$$F_\Phi = \frac{1}{\Phi} w_0 - \Phi^* \mathcal{W}\ , \quad F_{\bar{\Phi}} = \frac{1}{\bar{\Phi}} w_0 - \bar{\Phi}^* \mathcal{W}, \tag{2.6}$$

where we have defined for convenience $\sigma \equiv 1/(S+\bar{S})$. Since the supergravity framework in which we work is valid only up to the scale M, we will exclude from our analysis cases with the vevs of X, Φ, $\bar{\Phi}$ or $Y=\infty$. We will denote a field and its vacuum value (vev) by the same symbol, since there is no possibility of confusion. We can distinguish the following possibilities:

- $w \neq 0$: The constraints (1.3) allow us to write the dilaton F-term condition as $F_S = w(-r) = 0$, which requires $r=0$. But this is not possible, since we saw that r is a strictly positive quantity.

- $w = 0$.

  1. $S \neq \infty$ $(\sigma \neq 0)$: w can be zero only if $X=0$ provided $p<0$. However, since p is a strictly positive quantity, this is not an allowed solution.

  2. $S = \infty$ $(\sigma = 0)$: For $p>0$, there is a supersymmetric vacuum configuration with $X \neq 0$ and $\Phi = \bar{\Phi} = 0$.

We conclude that there is no supersymmetric minimum with a vanishing tree level cosmological constant and a finite value for the dilaton and therefore a nonzero gauge coupling.

U1MSB model

The superpotential in this model is:

$$\mathcal{W} = w_0 + w + k = \lambda \Theta X^2 + cX^{-p}e^{-rS} + k. \tag{2.7}$$

With the Kähler potential (1.2) the F-terms are:

$$F_\Theta = \frac{w_0}{\Theta} - \Theta^* \mathcal{W}, \quad F_S = -rw - \sigma \mathcal{W}, \quad F_X = \frac{1}{X}(w_0 - pw) - X^* \mathcal{W}. \tag{2.8}$$

Again, we distinguish two possibilities:

- $w \neq 0$: Using $\mathcal{W}=0$ we can write $F_S = w(-r) = 0$ and since $r>0$, there is no such supersymmetric vacuum.

- w=0:

  1. $S \neq \infty$ $(\sigma \neq 0)$: w can be zero only if $\Theta = 0$, which requires $p<0$. Since p is strictly positive, this is not an allowed solution.

2. $S = \infty$ ($\sigma = 0$): The $F_\Theta = 0$ equation implies $X = 0$. There is a supersymmetric vacuum with $p > 0$ and Θ undetermined.

We conclude that in the U1MSB model too, there is no supersymmetric vacuum except with zero gauge coupling.

3 Supersymmetry Breaking Vacua

GMSB model

In the GMSB model, there are four types of fields, namely the hidden sector field X, the messengers Φ and $\bar{\Phi}$ and the dilaton S. The physically relevant case is when the standard model nonsinglet messenger fields do not take vacuum expectation values, i.e. $\Phi = \bar{\Phi} = 0$, which implies through equations (2.6) that $F_\Phi = F_{\bar{\Phi}} = 0$. The minimization conditions simplify considerably if we notice that the values of the fields X and σ that minimize the potential with $V = 0$, will minimize $\tilde{V} \equiv \frac{V}{3|\mathcal{W}|^2}$ as well, since

$$\tilde{V}^{(\phi_i)} = \frac{V^{(\phi_i)}}{3|\mathcal{W}|^2} - \frac{V}{3|\mathcal{W}|^4}(|\mathcal{W}|^2)^{(\phi_i)}, \tag{3.9}$$

and that w_0 is essentially absent from the picture due to the vanishing of Φ and $\bar{\Phi}$. Define then

$$\tilde{F}_a \equiv \frac{F_X}{\sqrt{3}\mathcal{W}} = \frac{-1}{\sqrt{3}}\Big(\frac{p(1-\tilde{k})}{X} + X^*\Big) = \frac{-1}{\sqrt{3}}\Big(\frac{p(1-\tilde{k})}{a} + a\Big)e^{-i\alpha}, \tag{3.10}$$

where $\tilde{k} \equiv k/\mathcal{W}$. We derived the second part of the above equation using the definition of the F-term and the expression for $\mathcal{W}$ and in the third part of the equation we have defined $X \equiv ae^{i\alpha}$. Similarly, we define

$$\tilde{F}_S \equiv \frac{F_S}{\sqrt{3}\mathcal{W}} = \frac{-1}{\sqrt{3}}\Big(r(1-\tilde{k}) + \sigma\Big). \tag{3.11}$$

The phase in the above is zero if $k = 0$. These definitions allow us to write $\tilde{V} = |\tilde{F}_a|^2 + |\tilde{F}_S|^2 - 1$. Now let us look for solutions to the minimization conditions with $V = 0$ and try to find out if it is possible to generate the scale hierarchy necessary for a low energy gauge mediated supersymmetry breaking scenario. For simplicity, we will ignore all phases. The vacuum conditions become:

$$\tilde{V}^{(a)}: \quad \tilde{F}_a\Big[1-\frac{p}{a^2}(1-\tilde{k})(1+p\tilde{k})\Big]+\tilde{F}_S\Big[-\frac{pr}{a}\tilde{k}(1-\tilde{k})\Big]=0, \tag{3.12}$$

$$\tilde{V}^{(S)}: \quad \tilde{F}_a\Big[-\frac{pr}{a}\tilde{k}(1-\tilde{k})\Big]+\tilde{F}_S\Big[-r^2\tilde{k}(1-\tilde{k})-\sigma^2\Big]=0, \tag{3.13}$$

$$\tilde{V}=0: \quad \Big[\frac{p}{a}(1-\tilde{k})+a\Big]^2+\Big[r(1-\tilde{k})+\sigma\Big]^2=3. \tag{3.14}$$

We can rewrite (3.12) and (3.13) as

$$\frac{\tilde{F}_a}{\tilde{F}_S}=-\frac{ar}{p}-\frac{a\sigma^2}{pr\tilde{k}(1-\tilde{k})}=\frac{\frac{p}{a}(1-\tilde{k})+a}{r(1-\tilde{k})+\sigma}, \tag{3.15}$$

$$\Big[1-\frac{p}{a^2}(1-\tilde{k})(1+p\tilde{k})\Big]\Big[-r^2\tilde{k}(1-\tilde{k})-\sigma^2\Big]+\Big[-\frac{pr}{a}\tilde{k}(1-\tilde{k})\Big]\Big[-\frac{pr}{a}\tilde{k}(1-\tilde{k})\Big]=0. \tag{3.16}$$

The first part of the equation (3.15) comes from the vacuum condition $\tilde{V}^{(S)}=0$ and the second from the expressions (3.10) and (3.11) for the F-terms. Having in mind the constraint $a\leq\mathcal{O}(\eta)\sim 10^{-6}$ mentioned in the introduction and that $r\sim\mathcal{O}(10)$, $p\sim\mathcal{O}(1)$, we first notice that (3.14) can be satisfied only if $(1-\tilde{k})$ is also small, say $(1-\tilde{k})\sim\mathcal{O}(\eta\epsilon)$, with ϵ at most of order one. Then, (3.15) can be written as

$$\mathcal{O}\Big(\eta r+\frac{\sigma^2}{r\epsilon}\Big)=\mathcal{O}\Big(\frac{\epsilon+\eta}{r\eta\epsilon+\sigma}\Big). \tag{3.17}$$

We distinguish two possibilities. First, assume that $\sigma\sim\mathcal{O}(1)$. Then, (3.17) can be solved if $\epsilon\sim\mathcal{O}(\frac{1}{\sqrt{r}})$. Using, however, this value for ϵ, (3.16) becomes $\mathcal{O}(\frac{1}{\eta\sqrt{r}})\sim\mathcal{O}(r)$, which can not be satisfied since $\eta<<1$. If on the other hand, $\sigma<\mathcal{O}(1)$, then from (3.14) we see that $\epsilon\sim\mathcal{O}(1)$ and (3.15) becomes

$$\mathcal{O}\Big(\eta r+\frac{\sigma^2}{r}\Big)=\mathcal{O}\Big(\frac{1}{r\eta+\sigma}\Big). \tag{3.18}$$

We now turn to (3.16), which can be solved only if $\sigma^2\sim\mathcal{O}(\eta r^2)$. Substituting this back into (3.18), we get the condition $\mathcal{O}(\eta r)\sim\mathcal{O}(\frac{1}{r\sqrt{\eta}})$, which is not satisfied for $\eta<<1$. [4] We conclude that we can not satisfy the vacuum equations with $a\sim\mathcal{O}(\eta)$ and a vanishing tree level cosmological constant.

[4] If we relax the constraint $a\sim\eta$, we can solve the vacuum equations with $\sigma\sim 1$, and a, $(1-\tilde{k})\sim 1/r$. This seems to be the only possibility to avoid the negative conclusions of our analysis.

It is natural to ask how general this conclusion is, given that we have considered only a certain class of superpotentials. More specifically, we have not considered additional fields and couplings possible in the superpotential involving the field X. Such a coupling, can be of the form of either XAB or XXA or XXX, where A and B are chiral superfields of the hidden sector. Clearly, the most important contribution comes from XAB, when $A, B \sim \mathcal{O}(1)$. Then, defining $l \equiv < AB >$, $\tilde{l} \equiv l/\mathcal{W}$ and ignoring phases, we get the leading order modification to the F_a-term:

$$\sqrt{3}\tilde{F}_a \rightarrow \sqrt{3}\tilde{F}_a + \frac{p}{a}\tilde{l}, \tag{3.19}$$

which amounts to $\tilde{k} \rightarrow \tilde{k}'$, with $\tilde{k}' = \tilde{k} + \tilde{l}$. By similar arguments as before, we can see that our previous conclusion remains.

U1MSB model

The first question to address is if the $U(1)$ symmetry is flavor universal in the visible sector or not. The main argument for the existence of the pseudoanomalous $U(1)$ is that, provided that it is flavor dependent, it is an excellent candidate to explain fermion mass hierarchies [8]. But then, after supersymmetry breaking, its D-term will contribute to the soft masses which in turn tend to give large flavor changing neutral current (fcnc) contributions. Without any uneven mass splittings between squark generations, the U1MSB scenario is therefore consistent with data only if the dilaton F-term dominates over the D-terms. Here, we will assume that the D-term is negligible (under some additional assumptions it can indeed be small [9]) and argue that there exists a supersymmetry breaking vacuum with vanishing tree level cosmological constant.

The most important feature in the U1MSB model is that we do not have to make any special assumptions about the vevs of any of the fields. Each minimization condition can in principle determine its corresponding field vev. Indeed, a solution to the minimization conditions, as an expansion in the small parameter $\epsilon \equiv \Theta^2/Y^2$, has been presented in [5]. Finally, cancelation of the tree level cosmological constant can be achieved by an appropriate choice of k. Therefore, in this model, the vacuum conditions can be satisfied and at the same time a supersymmetry breaking scale of $\mathcal{O}(100)$ GeV can be generated. This scenario of supersymmetry breaking, in the dilaton dominance limit, qualitatively is very similar to gravity mediation. The difference comes from the field Θ which when taking a vev, provides us with an additional supersymmetry breaking parameter.

4 One loop cosmological constant

To ensure the vanishing of the cosmological constant at one loop, one has to look at the one loop effective potential [10]:

$$V_1 = V + \frac{1}{64\pi^2} Str\mathcal{M}^0 \cdot M^4 \log\frac{M^2}{\mu^2} + \frac{1}{32\pi^2} Str\mathcal{M}^2 \cdot M^2 + \frac{1}{64\pi^2} Str\mathcal{M}^4 \log\frac{\mathcal{M}^2}{M^2}, \tag{4.20}$$

with

$$Str\mathcal{M}^n = \sum_i (-1)^{2J_i}(2J_i+1)m_i^n \tag{4.21}$$

the well known supertrace formula of supergravity. Clearly, to have a zero cosmological constant, it is not sufficient to set $V=0$ alone. The second term is zero in theories with equal number of bosons and fermions and in particular in all supersymmetric models, but the third and the last term have to be taken into account. The last term is one that does not destabilize the hierarchy, but its presence is important for the vanishing of the cosmological constant. It can be taken into account by modifying the constraint on the classical potential $V=0$ to $V + 1/(64\pi^2)Str\mathcal{M}^4 \log\frac{\mathcal{M}^2}{M^2} = 0$. The additional term, however, is expected to be negligible compared to V, so we could safely ignore it in the previous section. In fact, we used the constant k introduced in the superpotential to carry out this "continuous" ($V=0$) fine tuning. The vanishing of the third term on the other hand is required in order to have a stable hierarchy, and it involves a "discrete" fine tuning. Given that $Str\mathcal{M}^2 = 2Qm_{3/2}^2$, with $m_{3/2}$ the gravitino mass and

$$Q = N - 1 - \mathcal{G}^i H_{i\bar{j}}\mathcal{G}^{\bar{j}}, \tag{4.22}$$

with $H_{i\bar{j}} = \partial_i\partial_{\bar{j}} \log\det\mathcal{G}_{m\bar{n}} - \partial_i\partial_{\bar{j}}\log\det Re[f_{ab}]$, a necessary condition for the vanishing of the cosmological constant is $Q=0$ [11]. In the above, $\mathcal{G}_i = \frac{\partial\mathcal{G}}{\partial\phi^i}$, N is the total number of chiral multiplets, $f_{ab} = S\delta_{ab}$ is the gauge kinetic function and $\mathcal{G}_{m\bar{n}}$ is the Kähler metric of the Kähler manifold for the N chiral superfields. We can readily calculate Q for both cases. $\mathcal{G}_S = \frac{F_S}{\mathcal{W}}$, $\mathcal{G}_{\bar{S}} = \bar{\mathcal{G}}_S$ and the Kähler metric is $\mathcal{G}_{i\bar{j}} = diag(1/(S+\bar{S})^2, -1)$, where the -1 is to be understood as multiplied by the $(N-1)\times(N-1)$ unit matrix, where N is the total number of chiral superfields. This is true only if the only modulus in the model is the dilaton and the Kähler potential is as we have chosen it. In more complicated situations the above formulas have to be modified accordingly. Also, $\mathcal{G}^{i\bar{j}} = diag((S+\bar{S})^2, -1)$ and $\mathcal{G}^i = \mathcal{G}^{i\bar{j}}\mathcal{G}_{\bar{j}}$, $\mathcal{G}^{\bar{i}} = \mathcal{G}_j\mathcal{G}^{j\bar{i}}$. A simple computation then yields:

$$Q = N - 1 - 3(S+\bar{S})^2 |\frac{F_S}{\mathcal{W}}|^2. \tag{4.23}$$

Remembering that $S + \bar{S} = \frac{2}{g^2} \sim \mathcal{O}(1)$ and that in realistic models $N \sim \mathcal{O}(150)$ (in the second reference in [9] $N = 143$), we conclude that if $|\frac{F_S}{\mathcal{W}}| \sim \mathcal{O}(1)$ (this does not exclude the case $|\frac{F_X}{\mathcal{W}}| \sim \mathcal{O}(1)$), using the condition for the vanishing of the tree level cosmological constant $|\frac{F_S}{\mathcal{W}}| \simeq 3$, we get the following condition for $Q = 0$: $N - 1 \simeq \frac{36}{g^4}$ [5]. It is amusing to notice that $N = 150$ implies $\alpha \simeq 1/26$. In other words, in dilaton dominated models with $V = 0$ at the minimum (the same goes through for other moduli dominated models), the one loop cosmological constant vanishes for values of the coupling constant remarkably close to its unification value. The dilaton F-term, therefore, has to be of at least equal order of magnitude as the F_X term.

5 Conclusions

We analyzed two simple models of gauge mediated supersymmetry breaking in the context of supergravity and in particular we looked for supersymmetry breaking vacua with zero cosmological constant. To be able to cancel the tree level cosmolological constant, we allowed for a constant k in the superpotential. With a Kähler potential of the form (1.2) and $k = 0$ it is not only impossible to do the latter but also it does not seem possible to get a weak coupling minimum [12].

[5] This relation is rather stable. By including in K the $-3\log(T+\bar{T})$ term, it changes to $N \simeq \frac{36}{g^4}$.

For the GMSB model, we showed that it is not possible to have a supersymmetry breaking vacuum state with zero tree level cosmological constant, with X in the desired range for low energy gauge mediation. We can speculate what would happen if we relaxed some of our simplifying assumptions. We could generalize the Kähler potential by including T or U moduli but our conclusions regarding the vanishing of the cosmological constant are unlikely to be changed, except that some other modulus F-term or a linear combination of them would take the place of F_S. Finally, corrections to the tree level Kähler potential would be small in the weak coupling limit and thus unlikely to change these conclusions.

For the U1MSB model, on the other hand, we had to assume that the D-terms were negligible in order to avoid conflict with fcnc data. Provided that this was the case, we argued that there exist (dilaton dominated) supersymmetry breaking minima with vanishing tree level and one loop cosmological constant.

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